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Virtual Sfermion Effects on Vector-Boson Pair Production at e^+e^- Colliders

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Abstract

We study the quantum effects on vector-boson pair production in e^+e^- annihilation induced by the sleptons and squarks of the minimal supersymmetric extension of the standard model (MSSM) in the one-loop approximation. We list full analytic results, and quantitatively analyze the resulting deviation from the standard-model prediction of $e^+e^- \rightarrow W^+W^-$ for the supergravity-inspired MSSM. The latter can be rendered small throughout the whole parameter space by an appropriate choice of renormalization scheme.

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The production of W -boson pairs in e^+e^- annihilation offers a unique opportunity to probe the nonabelian gauge structure of the standard model (SM) at the tree level, which manifests itself in a distinctive cancellation between the s -, t -, and u -channel scattering amplitudes. This process is being studied experimentally with high precision at the CERN Large Electron-Positron Collider (LEP2) in the centre-of-mass (CM) energy range $2M_W \lesssim \sqrt{s} \lesssim 205$ GeV. At the same time, the related reactions $e^+e^- \rightarrow \gamma\gamma, \gamma Z, ZZ$ are being measured there, too. With a future e^+e^- linear supercollider, such as JLC, NLC, or TESLA, these measurements can be extended to higher energies, way up to the TeV range, and rendered more precise. On the theoretical side, enormous effort has been invested into the computation of the one-loop radiative corrections to the cross sections of these processes in the SM, both for on- [1,2] and off-shell [3] vector bosons, and useful low- [4] and high-energy [5] approximations have been elaborated; for a comprehensive review, see Ref. [6].

Significant deviations of the measured cross sections from their SM predictions could signal physics beyond the SM. Since the gauge couplings of the electron are so tightly constrained by low-energy and LEP1 data, such deviations should mainly originate from the triple gauge-boson couplings (TGC's). Generic parameterizations of the so-called anomalous TGC's were introduced and applied to the processes $e^+e^- \rightarrow W^+W^-, ZZ$ in Ref. [7]. In order to explain the physical origin of anomalous TGC's, it is necessary to consider specific new-physics scenarios. From the theoretical point of view, renormalizable extensions of the SM are most satisfactory. As a rule, the deviations are then induced through loop effects of new particles, which affect not only the TGC's, but also the vector-boson propagators and the renormalizations of the parameters and wave functions of the tree-level amplitudes. Such deviations were investigated at the one-loop level in Refs. [8,9] for models with a modified lepton sector including Majorana neutrinos and were found to be generally small. In Ref. [10], the one-loop radiative corrections to the anomalous parameters $\Delta\kappa_V$ and λ_V of the VW^+W^- TGC's, with $V = \gamma, Z$, were studied in the minimal supersymmetric extension of the SM (MSSM), using the pinch technique to render them gauge independent. The sfermion contributions were found to generally dominate the Higgs and gaugino contributions. This may be understood by observing that mass splittings between the up and down components of the sfermion doublets give rise to significant contributions [10] and that the sfermions come in large numbers, due to their multiplicities in flavour and colour.

In this letter, we calculate the sfermion contributions to the cross sections of $e^+e^- \rightarrow V_1V_2$, with $V_1V_2 = \gamma\gamma, \gamma Z, ZZ, W^+W^-$, at one loop in the MSSM. Preliminary results of this study were published in Ref. [11]. Our calculation proceeds along the lines of Ref. [9], which gives full analytic results. We use the conventions of Ref. [9] and list only those formulas which need to be substituted therein. In a way, our analysis extends Ref. [12], where the sfermion-induced radiative corrections to the processes $e^+e^- \rightarrow Zh^0$ and $Z \rightarrow \gamma h^0$, with h^0 being the lightest CP-even Higgs boson, were calculated at one loop in the MSSM. The authors of Ref. [10] did not list analytic results that could be compared with ours. In Refs. [13,14,15], the one-loop radiative correction to $e^+e^- \rightarrow W^+W^-$ in the MSSM was considered for the full supersymmetric particle spectrum, under the simplistic

assumption that the mass matrix of each sfermion flavour is proportional to the unit matrix, so that the two weak eigenstates are mass eigenstates with a common mass. We shall compare the sfermion loop correction of Refs. [13,14,15] with our result below. As we shall see later, the size of the correction is significantly affected by the sfermion mass splittings.

The Higgs sector of the MSSM is made up by two complex Higgs isodoublet of opposite hypercharge and accommodates five physical Higgs bosons: the neutral CP-even h^0 and H^0 bosons, the neutral CP-odd A^0 boson, and the charged H^\pm -boson pair. At the tree level, it has two free parameters, which are usually taken to be the mass m_A of the A^0 boson and the ratio $\tan\beta = v_2/v_1$ of the vacuum expectation values of the two Higgs doublets. For each of these Higgs bosons and each SM fermion and gauge boson there is a supersymmetric partner. Thus, the spectrum of states is more than doubled if one passes from the SM to the MSSM, which gives rise to a proliferation of parameters and weakens the predictive power of the theory. A canonical method to reduce the number of parameters is to embed the MSSM into a grand unified theory (GUT), e.g., a suitable supergravity (SUGRA) model, in such a way that it is recovered in the low-energy limit. The MSSM thus constrained is described by the following parameters at the GUT scale, which come in addition to $\tan\beta$ and m_A : the universal scalar mass m_0 , the universal gaugino mass $m_{1/2}$, the trilinear Higgs-sfermion coupling A , the bilinear Higgs coupling B , and the Higgs-higgsino mass parameter μ . Notice that m_A is then not an independent parameter anymore, but it is fixed through the renormalization group equation. The number of parameters can be further reduced by making additional assumptions. Unification of the tau and bottom Yukawa couplings at the GUT scale leads to a correlation between m_t and $\tan\beta$. Furthermore, if the electroweak symmetry is broken radiatively, then B and μ are determined up to the sign of μ . Finally, it turns out that the MSSM parameters are nearly independent of the value of A , as long as $|A| \lesssim 500$ GeV at the GUT scale.

We now present our analytic results. We denote the four-momenta of e^+ , e^- , and the two produced vector bosons, V_1 and V_2 , by p_+ , p_- , k_1 , and k_2 , and define the Mandelstam variables as $s = (p_+ + p_-)^2$, $t = (p_+ - k_1)^2$, and $u = (p_+ - k_2)^2$. Neglecting the electron mass, we have $s + t + u = M_1^2 + M_2^2$, where M_1 and M_2 are the masses of V_1 and V_2 , respectively. In this limit, also the s -channel contributions due to Higgs-boson exchanges vanish. Because each of the four processes $e^+e^- \rightarrow V_1V_2$ has more than one tree-level diagram, it is convenient to introduce helicity amplitudes $\mathcal{M}^\kappa(\lambda_1, \lambda_2, s, t)$, where κ , λ_1 , and λ_2 denote the helicities of e^- , V_1 , and V_2 in the CM frame, respectively. The e^+ helicity is then $-\kappa$. The helicity amplitudes \mathcal{M}^κ can be decomposed into the standard matrix elements \mathcal{M}_i^κ ($i = 0, \dots, 9$) [2], which are written down in Appendix B of Ref. [9]. In addition to those for $i = 0, \dots, 3$, which already appear at the tree level, we only need \mathcal{M}_9^κ for the present analysis.

The tree-level cross sections of $e^+e^- \rightarrow V_1V_2$ are well known and may be found in Eq. (3.1) of Ref. [9]. The sfermion-induced one-loop corrections receive contributions from diagrams containing self-energy corrections, vertex corrections, and counterterm insertions. We work in the Fermi-constant (G_F) formulation of the electroweak on-shell renormalization scheme, which is explained in the context of Eq. (4.1) in Ref. [9]. Specif-

ically, starting from the results in the pure on-shell renormalization scheme, which uses Sommerfeld's fine-structure constant α and the physical particle masses as basic parameters, we fix $\alpha = \sqrt{2}G_F \sin^2 \theta_w M_W^2 / \pi$, where θ_w is the weak mixing angle, and supplement the radiative corrections with the term $-2\Delta r$, where Δr [16] contains those radiative corrections to the muon lifetime which the SM or its extensions introduce on top of the purely photonic corrections from within the Fermi model. The sfermion contribution to Δr in the MSSM was examined in Ref. [12]. All the formulas listed in Section III and Appendix D of Ref. [9] carry over to the sfermion case, except for Eqs. (3.5) and (3.8), which give the transverse parts of the vector-boson vacuum polarizations $\Pi_T^{V_1 V_1}$ and the proper vertex corrections $\delta\mathcal{M}_V^\kappa$, respectively. The relevant Feynman rules for the MSSM sfermion sector are summarized in Appendix A of Ref. [12]. For each fermion flavour $Q = U, D$, where $U = \nu_e, \nu_\mu, \nu_\tau, u, c, t$ and $D = e, \mu, \tau, d, s, b$, there is a corresponding sfermion flavour, denoted by a tilde. Except for the sneutrinos, which we assume to be left handed, \tilde{Q} comes in two mass eigenstates $a = 1, 2$. The masses $M_{\tilde{Q}a}$ of the sfermions and their trilinear and quartic couplings, $\tilde{V}_{Q_a Q_b}'^{V_i}$ and $\tilde{U}_{Q_a Q_b}'^{V_i V_j}$, respectively, to the vector bosons $V_i = \gamma, Z, W$ are defined in Appendix A of Ref. [12]. In the absence of flavour-changing neutral currents, we have $Q = Q'$ in $\tilde{V}_{Q_a Q_b}'^{V_i}$ and $\tilde{U}_{Q_a Q_b}'^{V_i V_j}$ if $V_i, V_j = \gamma, Z$, which explains the notation $\tilde{V}_{Qab}^{V_i}$ and $\tilde{U}_{Qab}^{V_i V_j}$ used in Ref. [12]. As in Ref. [12], we neglect the Cabibbo-Kobayashi-Maskawa mixing, so that we may write \tilde{V}_{UaDb}^W and \tilde{U}_{Qab}^{WW} . The sfermion contributions to the $\Pi_T^{V_1 V_2}$ functions read [12]

$$\begin{aligned} \Pi_T^{V_1 V_2}(p^2) &= \frac{1}{48\pi^2} \sum_{Q,a,b} N_{\text{col}}^Q \tilde{V}_{Q_a Q_b}^{V_1} \tilde{V}_{Q_b Q_a}^{V_2} \left\{ \left[s - 2(M_{\tilde{Q}_a}^2 + M_{\tilde{Q}_b}^2) + \frac{(M_{\tilde{Q}_a}^2 - M_{\tilde{Q}_b}^2)^2}{s} \right] \right. \\ &\quad \times B_0(p, M_{\tilde{Q}_a}, M_{\tilde{Q}_b}) + M_{\tilde{Q}_a}^2 \left(2 - \frac{M_{\tilde{Q}_a}^2 - M_{\tilde{Q}_b}^2}{s} \right) B_0(0, M_{\tilde{Q}_a}, M_{\tilde{Q}_a}) \\ &\quad \left. + M_{\tilde{Q}_b}^2 \left(2 - \frac{M_{\tilde{Q}_b}^2 - M_{\tilde{Q}_a}^2}{s} \right) B_0(0, M_{\tilde{Q}_b}, M_{\tilde{Q}_b}) + \frac{2}{3}s - \frac{(M_{\tilde{Q}_a}^2 - M_{\tilde{Q}_b}^2)^2}{s} \right\}, \quad (1) \end{aligned}$$

where p is the external four-momentum, $N_{\text{col}}^Q = 1$ (3) for sleptons (squarks), and the standard two-point scalar function B_0 is defined in Eq. (C.2) of Ref. [9]. If $V_1 = W^-$ and $V_2 = W^+$, then $Q_a = U_a$, $Q_b = D_b$, and it is summed over (U, D) instead of Q . The sfermion contribution to $\delta\mathcal{M}_V^\kappa$ is found to be

$$\begin{aligned} \delta\mathcal{M}_V^\kappa &= -\frac{1}{2\pi^2} \sum_{B=\gamma,Z} \frac{g_{eeB}^\kappa}{s - M_B^2} \sum_{Q,a,b,c} N_{\text{col}}^Q (\tilde{V}_{Q_c Q_b}^B \tilde{V}_{Q_b Q_a}^{V_1} \tilde{V}_{Q_a Q_c}^{V_2} - \tilde{V}_{Q_c Q_a}^{V_2} \tilde{V}_{Q_a Q_b}^{V_1} \tilde{V}_{Q_b Q_c}^B) \\ &\quad \times [\mathcal{M}_1^\kappa (C_2^0 + C_3^{01} + C_3^{02}) + \mathcal{M}_2^\kappa C_3^{01} + \mathcal{M}_3^\kappa C_3^{02} - \mathcal{M}_9^\kappa (C_2^{12} + C_3^{12} + C_3^{21})], \quad (2) \end{aligned}$$

where the C functions are the Lorentz coefficients of the standard three-point tensor integrals defined in Eq. (C4) of Ref. [9]. In Eq. (2), we have suppressed their common

argument $(k_1, -k_2, M_{\tilde{Q}_a}, M_{\tilde{Q}_b}, M_{\tilde{Q}_c})$. Deviating from Ref. [9], the electron gauge couplings appearing in Eq. (2) are defined as

$$g_{ee\gamma}^{\pm} = e, \quad g_{eeZ}^{+} = -\frac{e \sin \theta_w}{\cos \theta_w}, \quad g_{eeZ}^{-} = \frac{e}{\cos \theta_w \sin \theta_w} \left(\frac{1}{2} - \sin^2 \theta_w \right), \quad (3)$$

where $e = \sqrt{4\pi\alpha}$. We caution the reader that the Feynman rules used in Refs. [9,12] differ in the sign of $\sin \theta_w$. Consequently, we need to multiply the expression for $\Pi_T^{\gamma Z}$ in Eq. (1) with an extra minus sign when we insert it into the relevant formulas, Eqs. (3.6) and (D2), of Ref. [9]. Notice that Eq. (2) includes the contributions from both the direct and crossed triangle diagrams, which are proportional to $\tilde{V}_{Q_c Q_b}^B \tilde{V}_{Q_b Q_a}^{V_1} \tilde{V}_{Q_a Q_c}^{V_2}$ and $\tilde{V}_{Q_c Q_a}^{V_2} \tilde{V}_{Q_a Q_b}^{V_1} \tilde{V}_{Q_b Q_c}^B$, respectively. If $V_1 = W^-$ and $V_2 = W^+$, then the first term contributes for $Q_a = U_a$, $Q_b = D_b$, and $Q_c = D_c$ and the second one for $Q_a = D_a$, $Q_b = U_b$, and $Q_c = U_c$.

At this point, we should compare our results with those published in Refs. [13,14,15]. To that end, we put $M_{\tilde{Q}_1} = M_{\tilde{Q}_2}$ and nullify the mixing angle relating the weak and mass eigenstates for each sfermion flavour \tilde{Q} . Then, our Eq. (1) agrees with Eqs. (C7), (D5), (E9), and (F8) in Ref. [13], up to an overall minus sign, if we eliminate the factor $1/2$ multiplying T_{3f}^i in Eq. (D5) and the sum over i in Eq. (F8). As for the γW^+W^- vertex correction, our Eq. (2) is in accordance with Eqs. (65) and (B13)–(B18) in Ref. [14] if we replace the first two appearances of C_{36} in Eq. (B14) by C_{35} , substitute C_{24} in Eq. (B18) by C_{22} , and include an overall minus sign in Eqs. (B17) and (B18). As for the $Z W^+W^-$ vertex correction, we find agreement with Eqs. (97) and (C48)–(C53) in Ref. [14] if we alter the overall signs of Eqs. (C52) and (C53).

Now, we explore the phenomenological implications of our results. We concentrate on the case of $e^+e^- \rightarrow W^+W^-$ because, for $\sqrt{s} \gtrsim 180$ GeV, it has the largest cross section of the four processes under consideration and it is the only one involving TGC's at the tree level in the SM. The SM input parameters for our numerical analysis are taken to be $G_F = 1.16639 \cdot 10^{-5}$ GeV $^{-2}$ [17], $m_W = 80.385$ GeV, $m_Z = 91.1871$ GeV, $m_t = 174.3$ GeV [18], and $m_b = 4.7$ GeV. We vary $\tan \beta$ and m_A in the ranges $1 < \tan \beta < 35 \approx m_t/m_b$ and 100 GeV $< m_A < 600$ TeV, respectively. As for the GUT parameters, we choose $m_{1/2} = 150$ GeV, $A = 0$, and $\mu < 0$, and tune m_0 so as to be consistent with the desired value of m_A . All other MSSM parameters are then determined according to the SUGRA-inspired scenario as implemented in the program package SUSPECT [19]. We checked that the results obtained from the program package ISAJET 7.49 [20], where the electroweak-symmetry-breaking scale is fixed to be $Q = \sqrt{M_{\tilde{t}_L} M_{\tilde{t}_R}}$, agree with those from SUSPECT within typically 5% or less if the same scale convention is implemented in the latter. In our analysis, we adopt the SUSPECT default value $Q = M_Z$. We do not impose the unification of the tau and bottom Yukawa couplings at the GUT scale, which would just constrain the allowed $\tan \beta$ range without any visible effect on the results for these values of $\tan \beta$. We exclude solutions which do not comply with the present experimental lower mass bounds of the sfermions, charginos, neutralinos, and Higgs bosons [21].

In Fig. 1, the sfermion-induced correction $\delta(\theta)$ in the relationship $d\sigma/d\cos\theta = (d\sigma/d\cos\theta)_{\text{Born}}[1 + \delta(\theta)]$ between the one-loop-corrected and tree-level cross sections of

$e^+e^- \rightarrow W^+W^-$ is shown as a function of the scattering angle θ , enclosed between the e^+ and W^+ three-momenta in the CM frame, for $\sqrt{s} = 200, 500$ and 1000 GeV assuming $\tan\beta = 10$ and $m_A = 250$ GeV. We observe that $\delta(\theta)$ has a typical size of order 0.1% or less and can be of either sign. In the backward direction, it strongly depends on the CM energy, while the energy dependence is rather feeble in the forward direction. We emphasize that the smallness of $\delta(\theta)$ is a special feature of the G_F scheme. In the α scheme, the correction is given by $\delta(\theta) + 2\Delta r$ and thus shifted to negative values because we have $\Delta r \approx -0.11\%$, as indicated in Fig. 1. Next, we study the correction Δ to the integrated cross section, defined by $\sigma = \sigma_{\text{Born}}(1 + \Delta)$, for $\sqrt{s} = 200$ GeV. In Fig. 2, the $\tan\beta$ dependences of Δ and Δr are shown for $m_A = 100, 250$, and 600 GeV, while, in Fig. 3, the m_A dependences are shown for $\tan\beta = 3, 10$, and 30 . These dependences are implicit in the sense that our formulas for Δ and Δr do not contain $\tan\beta$ or m_A . In fact, Δ and Δr only depend on $\tan\beta$ or m_A via the sfermion masses and gauge couplings, the latter being affected through the mixing angles which rotate the weak eigenstates of the sfermion into their mass eigenstates. We note that the SUGRA-inspired MSSM with our choice of input parameters does not permit $\tan\beta$ and m_A to be simultaneously small, due to the experimental selectron mass lower bound [21]. This explains why the curves for $m_A = 100$ GeV in Fig. 2 only start at $\tan\beta \approx 11$ and those for $\tan\beta = 3$ in Fig. 3 at $m_A \approx 240$ GeV. For large m_A , the experimental m_h lower bound [21] enforces $\tan\beta \gtrsim 3$. On the other hand, the experimental lower bounds on the chargino and neutralino masses [21] induce an upper limit on $\tan\beta$, which depends on m_A . From Fig. 2 we observe that the $\tan\beta$ dependence of Δr for fixed m_A is modest for intermediate values of $\tan\beta$, while Δr increases in magnitude towards the edges of the allowed $\tan\beta$ range. The stau and tau-sneutrino contributions dominate for large $\tan\beta$ and small m_A , while the sbottom and stop contributions dominate for small $\tan\beta$ and large m_A . The contributions due to the sfermions of the first and second generations are insignificant for all values of $\tan\beta$ and m_A . It is interesting to investigate the mixings between the left- and right-handed components of the charged sfermions in the third generation. The mixing is strongest for stop, especially for small $\tan\beta$ and small m_A . For stau and sbottom, the mixings are generally feeble for large $\tan\beta$, independently of m_A . The magnitude of Δr may reach several tenths of percent if m_A is small to medium and $\tan\beta$ is close to its lower or upper limits. For $\tan\beta \approx 30$, Δr is almost independent of m_A , while for smaller (larger) values of $\tan\beta$, the size of Δr monotonically decreases (increases) as m_A increases. These features are also nicely illustrated in Fig. 3. We learn from Figs. 2 and 3 that Δ is insignificant, below 0.02% in size, for all considered values of $\tan\beta$ and m_A . We stress that this happens by virtue of the G_F scheme.

In summary, we derived analytic results for the sfermion-induced radiative corrections to the cross sections of $e^+e^- \rightarrow \gamma\gamma, \gamma Z, ZZ, W^+W^-$ at one loop in the MSSM and presented a phenomenological discussion for the most interesting case, $e^+e^- \rightarrow W^+W^-$, adopting a SUGRA-inspired scenario. In the latter case, the correction can essentially be quenched by adopting the G_F scheme, which could not be anticipated without explicit calculation. On the other hand, the sfermions are likely to generate the bulk of the MSSM correction to $e^+e^- \rightarrow W^+W^-$ because of their multiplicities in flavour and

colour. This expectation is substantiated by a study of the MSSM corrections to the γW^+W^- and ZW^+W^- TGC's [10]. We conclude that significant deviations of the measured cross section of $e^+e^- \rightarrow W^+W^-$ from its SM predictions will not point towards the SUGRA-inspired MSSM.

Note added

After the completion of this work, we received a preprint [22] which reports on the MSSM sfermion corrections to the cross section of $e^+e^- \rightarrow W^+W^-$ in the modified minimal-subtraction scheme. The analytic results for the vector-boson vacuum polarizations and the proper vertex corrections given in Eqs. (B.1), (B.3)–(B.5), (B.8) and (B.9) of Ref. [22] agree with our Eqs. (1) and (2), respectively. The agreement was also established numerically to very high precision.

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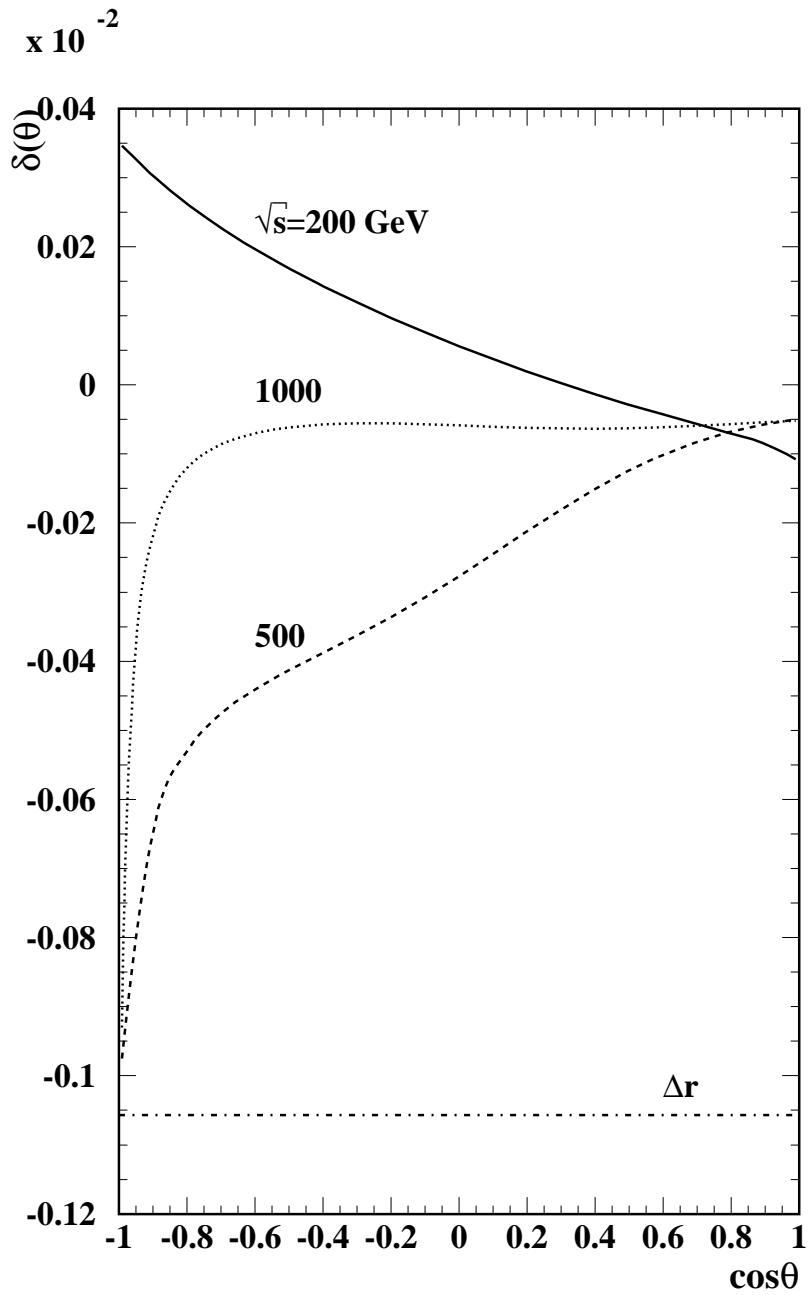


Figure 1: Relative correction $\delta(\theta)$ to the differential cross section $d\sigma/d\cos\theta$ of $e^+e^- \rightarrow W^+W^-$, for $\sqrt{s} = 200$, 500, and 1000 GeV, and contribution to Δr due to the sfermions in the SUGRA-inspired MSSM with $\tan\beta = 10$ and $m_A = 250$ GeV, as functions of $\cos\theta$.

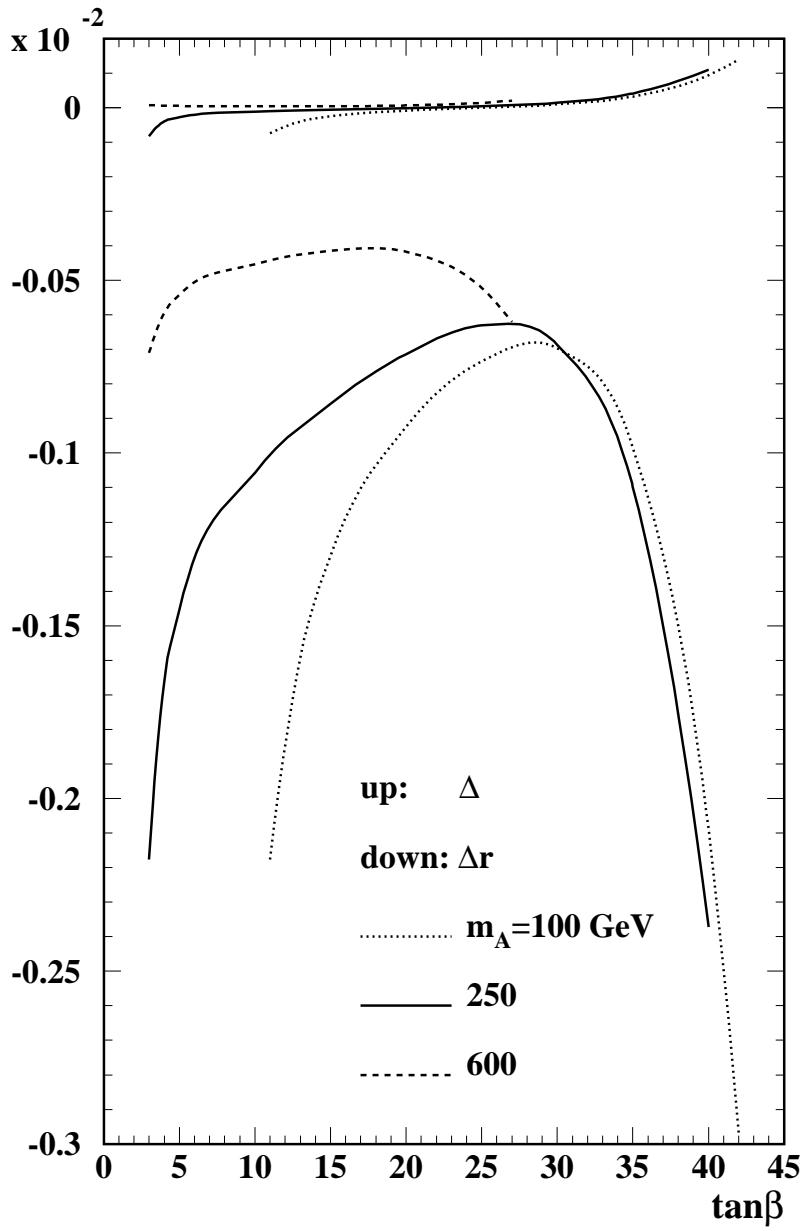


Figure 2: Relative correction Δ to the total cross section σ of $e^+e^- \rightarrow W^+W^-$, for $\sqrt{s} = 200 \text{ GeV}$, and contribution to Δr due to the sfermions in the SUGRA-inspired MSSM with $m_A = 100, 250$, and 600 GeV , as functions of $\tan\beta$.

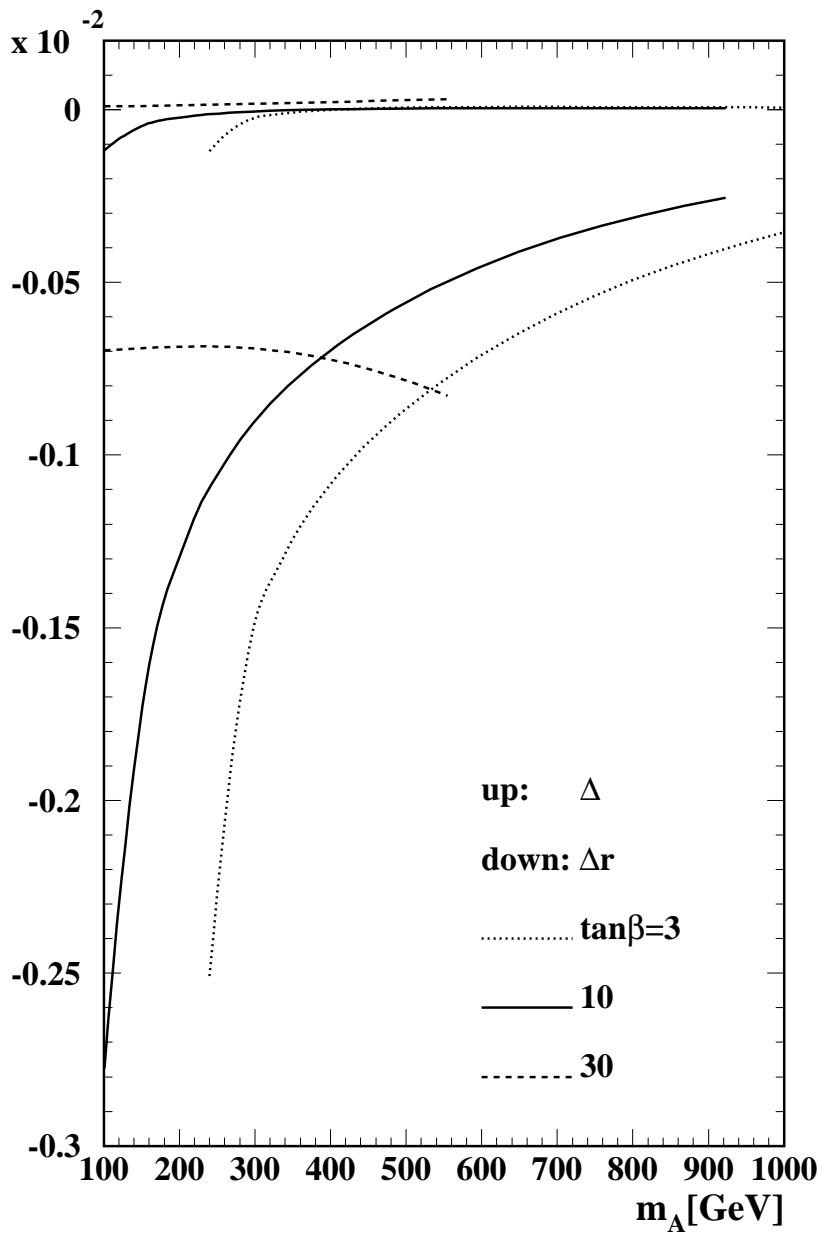


Figure 3: Relative correction Δ to the total cross section σ of $e^+e^- \rightarrow W^+W^-$, for $\sqrt{s} = 200$ GeV, and contribution to Δr due to the sfermions in the SUGRA-inspired MSSM with $\tan\beta = 3, 10$, and 30 , as functions of m_A .